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# Phenomenology of the utilitarian supersymmetric standard model

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**Abstract**

We study the 2010 specific version of the 2002 proposed $U(1)_X$ extension of the supersymmetric standard model, which has no $\mu$ term and conserves baryon number and lepton number separately and automatically. We consider in detail the scalar sector as well as the extra $Z_X$ gauge boson, and their interactions with the necessary extra color-triplet particles of this model, which behave as leptoquarks. We show how the diphoton excess at 750 GeV, recently observed at the LHC, may be explained within this context. We identify a new fermion dark-matter candidate and discuss its properties. An important byproduct of this study is the discovery of relaxed supersymmetric constraints on the Higgs boson's mass of 125 GeV.


## 1. Introduction

Since the recent announcements [1,2] by the ATLAS and CMS Collaborations at the Large Hadron Collider (LHC) of a diphoton excess around 750 GeV, numerous papers [3] have appeared explaining its presence or discussing its implications. In this paper, we study the phenomenology of a model proposed in 2002 [4], which has exactly all the necessary and sufficient particles and interactions for this purpose. They were of course there for solving other issues in

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Table 1
Particle content of proposed model.

| Superfield | $SU(3)_C$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)_X$ |
|---|---|---|---|---|
| $Q=(u,d)$ | 3 | 2 | 1/6 | 0 |
| $u^c$ | 3* | 1 | −2/3 | 1/2 |
| $d^c$ | 3* | 1 | 1/3 | 1/2 |
| $L=(\nu,e)$ | 1 | 2 | −1/2 | 1/3 |
| $e^c$ | 1 | 1 | 1 | 1/6 |
| $N^c$ | 1 | 1 | 0 | 1/6 |
| $\phi_1$ | 1 | 2 | −1/2 | −1/2 |
| $\phi_2$ | 1 | 2 | 1/2 | −1/2 |
| $S_1$ | 1 | 1 | 0 | −1/3 |
| $S_2$ | 1 | 1 | 0 | −2/3 |
| $S_3$ | 1 | 1 | 0 | 1 |
| $U$ | 3 | 1 | 2/3 | −2/3 |
| $D$ | 3 | 1 | −1/3 | −2/3 |
| $U^c$ | 3* | 1 | −2/3 | −1/3 |
| $D^c$ | 3* | 1 | 1/3 | −1/3 |

particle physics. However, the observed diphoton excess may well be a first revelation [5] of this model, including its connection to dark matter.

This 2002 model extends the supersymmetric standard model by a new $U(1)_X$ gauge symmetry. It replaces the $\mu$ term with a singlet scalar superfield which also couples to heavy color-triplet superfields which are electroweak singlets. The latter are not *ad hoc* inventions, but are necessary for the cancellation of axial-vector anomalies. It was shown in Ref. [4] how this was accomplished by the remarkable exact factorization of the sum of eleven cubic terms, resulting in two generic classes of solutions [6]. Both are able to enforce the conservation of baryon number and lepton number up to dimension-five terms. As such, the scalar singlet and the vectorlike quarks are indispensible ingredients of this 2002 model. They are thus naturally suited for explaining the observed diphoton excess. In 2010 [7], a specific version was discussed, which will be the subject of this paper as well. An important byproduct of this study is the discovery of relaxed supersymmetric constraints on the Higgs boson's mass of 125 GeV. This is independent of whether the diphoton excess is confirmed or not.

## 2. Model

Consider the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$ with the particle content of Ref. [4]. For $n_1 = 0$ and $n_4 = 1/3$ in Solution (A), the various superfields transform as shown in Table 1. There are three copies of $Q, u^c, d^c, L, e^c, N^c, S_1, S_2$; two copies of $U, U^c, S_3$; and one copy of $\phi_1, \phi_2, D, D^c$. The only allowed terms of the superpotential are thus trilinear, i.e.

$$Qu^c\phi_2, \quad Qd^c\phi_1, \quad Le^c\phi_1, \quad LN^c\phi_2, \quad S_3\phi_1\phi_2, \quad N^cN^cS_1, \tag{1}$$

$$S_3UU^c, \quad S_3DD^c, \quad u^cN^cU, \quad u^ce^cD, \quad d^cN^cD, \quad QLD^c, \quad S_1S_2S_3. \tag{2}$$

The absence of any bilinear term means that all masses come from soft supersymmetry breaking, thus explaining why the $U(1)_X$ and electroweak symmetry breaking scales are not far from that of supersymmetry breaking. As $S_{1,2,3}$ acquire nonzero vacuum expectation values (VEVs), the exotic $(U, U^c)$ and $(D, D^c)$ fermions obtain Dirac masses from $\langle S_3 \rangle$, which also generates the $\mu$

term. The singlet $N^c$ fermion gets a large Majorana mass from $\langle S_1 \rangle$, so that the neutrino $\nu$ gets a small seesaw mass in the usual way. The singlet $S_{1,2,3}$ fermions themselves get Majorana masses from their scalar counterparts $\langle S_{1,2,3} \rangle$ through the $S_1 S_2 S_3$ terms. The only massless fields left are the usual quarks and leptons. They then become massive as $\phi^0_{1,2}$ acquire VEVs, as in the minimal supersymmetric standard model (MSSM).

Because of $U(1)_X$, the structure of the superpotential conserves both $B$ and $(-1)^L$, with $B = 1/3$ for $Q, U, D$, and $B = -1/3$ for $u^c, d^c, U^c, D^c$; $(-1)^L$ odd for $L, e^c, N^c, U, U^c, D, D^c$, and even for all others. Hence the exotic $U, U^c, D, D^c$ scalars are leptoquarks and decay into ordinary quarks and leptons. The $R$ parity of the MSSM is defined here in the same way, i.e. $R \equiv (-)^{2j+3B+L}$, and is conserved. Note also that the quadrilinear terms $QQQL$ and $u^c u^c d^c e^c$ (allowed in the MSSM) as well as $u^c d^c d^c N^c$ are forbidden by $U(1)_X$. Proton decay is thus strongly suppressed. It may proceed through the quintilinear term $QQQLS_1$ as the $S_1$ fields acquire VEVs, but this is a dimension-six term in the effective Lagrangian, which is suppressed by two powers of a very large mass, say the Planck mass, and may safely be allowed.

## 3. Gauge sector

The new $Z_X$ gauge boson of this model becomes massive through $\langle S_{1,2,3} \rangle = u_{1,2,3}$, whereas $\langle \phi^0_{1,2} \rangle = v_{1,2}$ contribute to both $Z$ and $Z_X$. The resulting $2 \times 2$ mass-squared matrix is given by [8]

$$\mathcal{M}^2_{Z,Z_X} = \begin{pmatrix} (1/2)g_Z^2(v_1^2 + v_2^2) & (1/2)g_Z g_X (v_2^2 - v_1^2) \\ (1/2)g_Z g_X (v_2^2 - v_1^2) & 2g_X^2[(1/9)u_1^2 + (4/9)u_2^2 + u_3^2 + (1/4)(v_1^2 + v_2^2)] \end{pmatrix}. \tag{3}$$

Since precision electroweak measurements require $Z - Z_X$ mixing to be very small [9], $v_1 = v_2$, i.e. $\tan\beta = 1$, is preferred. With the 2012 discovery [10,11] of the 125 GeV particle, and identified as the one Higgs boson $h$ responsible for electroweak symmetry breaking, $\tan\beta = 1$ is not compatible with the MSSM, but is perfectly consistent here, as shown already in Ref. [7] and in more detail in the next section.

Consider the decay of $Z_X$ to the usual quarks and leptons. Each fermionic partial width is given by

$$\Gamma(Z_X \to \bar{f} f) = \frac{g_X^2 M_{Z_X}}{24\pi}[c_L^2 + c_R^2], \tag{4}$$

where $c_{L,R}$ can be read off under $U(1)_X$ from Table 1. Thus

$$\frac{\Gamma(Z_X \to \bar{t}t)}{\Gamma(Z_X \to \mu^+\mu^-)} = \frac{\Gamma(Z_X \to \bar{b}b)}{\Gamma(Z_X \to \mu^+\mu^-)} = \frac{27}{5}. \tag{5}$$

This will serve to distinguish it from other $Z'$ models [12].

At the LHC, limits on the mass of any $Z'$ boson depend on its production by $u$ and $d$ quarks times its branching fraction to $e^-e^+$ and $\mu^-\mu^+$. In a general analysis of $Z'$ couplings to $u$ and $d$ quarks,

$$\mathcal{L} = \frac{g'}{2} Z'_\mu \bar{f} \gamma_\mu (g_V - g_A \gamma_5) f, \tag{6}$$

where $f = u, d$. The $c_u, c_d$ coefficients used in an experimental search [13,14] of $Z'$ are then given by

$$c_u = \frac{g'^2}{2}[(g_V^u)^2 + (g_A^u)^2]B(Z' \to l^-l^+), \quad c_d = \frac{g'^2}{2}[(g_V^d)^2 + (g_A^d)^2]B(Z' \to l^-l^+), \tag{7}$$

where $l = e, \mu$. In this model

$$c_u = c_d = \frac{g_X^2}{4}B(Z' \to l^-l^+). \tag{8}$$

To estimate $B(Z' \to l^-l^+)$, we assume $Z_X$ decays to all SM quarks and leptons with effective zero mass, all the scalar leptons with effective mass of 500 GeV, all the scalar quarks with effective mass of 800 GeV, the exotic $U, D$ fermions with effective mass of 400 GeV (needed to explain the diphoton excess), and one pseudo-Dirac fermion from combining $\tilde{S}_{1,2}$ (the dark matter candidate to be discussed) with mass of 200 GeV. We find $B(Z' \to l^-l^+) = 0.04$, and for $g_X = 0.53$, a lower bound of 2.85 TeV on $m_{Z_X}$ is obtained from the LHC data based on the 7 and 8 TeV runs.

## 4. Scalar sector

Consider the scalar potential consisting of $\phi_{1,2}$ and $S_{1,2,3}$. Whereas there are 2 copies of $S_3$ and 3 copies each of $S_{1,2}$, we can choose one copy each to be the one with nonzero vacuum expectation value. We then assume that the superpotential linking them is given by

$$W = fS_3\phi_1\phi_2 + hS_3S_2S_1, \tag{9}$$

which is of course missing some terms. We have neglected them for simplicity. Its contribution to the scalar potential is

$$V_F = f^2(\Phi_1^\dagger\Phi_1 + \Phi_2^\dagger\Phi_2)S_3^*S_3 + h^2(S_1^*S_1 + S_2^*S_2)S_3^*S_3 + |f\Phi_1^\dagger\Phi_2 + hS_1S_2|^2, \tag{10}$$

where $\phi_1$ has been redefined to $\Phi_1 = (\phi_1^+, \phi_1^0)$. The gauge contribution is

$$\begin{aligned} V_D = & \frac{1}{8}g_2^2[(\Phi_1^\dagger\Phi_1)^2 + (\Phi_2^\dagger\Phi_2)^2 + 2(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) - 4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1)] \\ & + \frac{1}{8}g_1^2[-(\Phi_1^\dagger\Phi_1) + (\Phi_2^\dagger\Phi_2)]^2 \\ & + \frac{1}{2}g_X^2\left[-\frac{1}{2}\Phi_1^\dagger\Phi_1 - \frac{1}{2}\Phi_2^\dagger\Phi_2 - \frac{1}{3}S_1^*S_1 - \frac{2}{3}S_2^*S_2 + S_3^*S_3\right]^2. \end{aligned} \tag{11}$$

The soft supersymmetry-breaking terms are

$$\begin{aligned} V_{soft} = & \mu_1^2\Phi_1^\dagger\Phi_1 + \mu_2^2\Phi_2^\dagger\Phi_2 + m_3^2S_3^*S_3 + m_2^2S_2^*S_2 + m_1^2S_1^*S_1 \\ & + [m_{12}S_2^*S_1^2 + A_f fS_3\Phi_1^\dagger\Phi_2 + A_h hS_3S_2S_1 + H.c.]. \end{aligned} \tag{12}$$

In addition, there is an important one-loop contribution from the $t$ quark and its supersymmetric scalar partners:

$$V_t = \frac{1}{2}\lambda_2(\Phi_2^\dagger\Phi_2)^2, \tag{13}$$

where

$$\lambda_2 = \frac{6G_F^2m_t^4}{\pi^2}\ln\left(\frac{m_{\tilde{t}_1}m_{\tilde{t}_2}}{m_t^2}\right) \tag{14}$$

is the well-known correction which allows the Higgs mass to exceed $m_Z$.

Let $\langle\phi^0_{1,2}\rangle = v_{1,2}$ and $\langle S_{1,2,3}\rangle = u_{1,2,3}$, we study the conditions for obtaining a minimum of the scalar potential $V = V_F + V_D + V_{soft} + V_t$. We look for the solution $v_1 = v_2 = v$ which implies that

$$\mu_1^2 = \mu_2^2 + \lambda_2 v^2 \tag{15}$$

$$0 = \mu_1^2 + A_f f u_3 + f^2(u_3^2 + v^2) + \frac{1}{2}g_X^2\left(v^2 + \frac{1}{3}u_1^2 + \frac{2}{3}u_2^2 - u_3^2\right) + f h u_1 u_2. \tag{16}$$

We then require that this solution does not mix the $Re(\phi_{1,2})$ and $Re(S_{1,2,3})$ sectors. The additional conditions are

$$0 = A_f f + (2f^2 - g_X^2)u_3, \tag{17}$$

$$0 = \frac{1}{3}g_X^2 u_1 + f h u_2, \tag{18}$$

$$0 = \frac{2}{3}g_X^2 u_2 + f h u_1. \tag{19}$$

Hence

$$u_1 = \sqrt{2}u_2, \quad f h = \frac{-\sqrt{2}g_X^2}{3}. \tag{20}$$

The $2 \times 2$ mass-squared matrix spanning $[\sqrt{2}Re(\phi_1^0), \sqrt{2}Re(\phi_2^0)]$ is

$$\mathcal{M}_\phi^2 = \begin{pmatrix} \kappa + g_X^2 v^2/2 & -\kappa + g_X^2 v^2/2 + 2f^2 v^2 \\ -\kappa + g_X^2 v^2/2 + 2f^2 v^2 & \kappa + g_X^2 v^2/2 + 2\lambda_2 v^2 \end{pmatrix}, \tag{21}$$

where

$$\kappa = (2f^2 - g_X^2)u_3^2 + \frac{2}{3}g_X^2 u_2^2 + \frac{1}{2}(g_1^2 + g_2^2)v^2. \tag{22}$$

For $\lambda_2 v^2 << \kappa$, the Higgs boson $h \simeq Re(\phi_1^0 + \phi_2^0)$ has a mass given by

$$m_h^2 \simeq \left(g_X^2 + 2f^2 + \lambda_2\right)v^2, \tag{23}$$

whereas its heavy counterpart $H \simeq Re(-\phi_1^0 + \phi_2^0)$ has a mass given by

$$m_H^2 \simeq (4f^2 - 2g_X^2)u_3^2 + \frac{4}{3}g_X^2 u_2^2 + (g_1^2 + g_2^2 - 2f^2 + \lambda_2)v^2. \tag{24}$$

The conditions for obtaining the minimum of $V$ in the $S_{1,2,3}$ directions are

$$0 = m_3^2 + g_X^2 u_3^2 + \left(3h^2 - \frac{4}{3}g_X^2\right)u_2^2 + \frac{\sqrt{2}A_h h u_2^2}{u_3}, \tag{25}$$

$$0 = m_2^2 + 2m_{12}u_2 + \left(2h^2 + \frac{8}{9}g_X^2\right)u_2^2 + \left(h^2 - \frac{2}{3}g_X^2\right)u_3^2 + \sqrt{2}A_h h u_3, \tag{26}$$

$$0 = m_1^2 + 2m_{12}u_2 + \left(h^2 + \frac{4}{9}g_X^2\right)u_2^2 + \left(h^2 - \frac{1}{3}g_X^2\right)u_3^2 + \frac{1}{\sqrt{2}}A_h h u_3. \tag{27}$$

The $3 \times 3$ mass-squared matrix spanning $[\sqrt{2}Re(S_1), \sqrt{2}Re(S_2), \sqrt{2}Re(S_3)]$ is given by

$$m_{11}^2 = \frac{4}{9} g_X^2 u_2^2 - \frac{1}{\sqrt{2}} A_h h u_3 + \frac{1}{3} g_X^2 v^2, \quad m_{22}^2 = 2m_{11}^2 - 2m_{12}u_2, \tag{28}$$

$$m_{12}^2 = m_{21}^2 = 2\sqrt{2} m_{12} u_2 + A_h h u_3 + 2\sqrt{2}\left(h^2 + \frac{2}{9} g_X^2\right) u_2^2 - \frac{\sqrt{2}}{3} g_X^2 v^2, \tag{29}$$

$$m_{33}^2 = 2g_X^2 u_3^2 - \sqrt{2} A_h h u_2^2/u_3 + (2f^2 - g_X^2) v^2, \tag{30}$$

$$m_{13}^2 = m_{31}^2 = A_h h u_2 + 2\sqrt{2}\left(h^2 - \frac{1}{3} g_X^2\right) u_3 u_2, \tag{31}$$

$$m_{23}^2 = m_{32}^2 = \sqrt{2} A_h h u_2 + 2\left(h^2 - \frac{2}{3} g_X^2\right) u_3 u_2. \tag{32}$$

The $5 \times 5$ mass-squared matrix spanning $[\sqrt{2}Im(\phi_1^0), \sqrt{2}Im(\phi_2^0), \sqrt{2}Im(S_1), \sqrt{2}Im(S_2), \sqrt{2}Im(S_3)]$ has two zero eigenvalues, corresponding to the would-be Goldstone modes

$$(1, 1, 0, 0, 0) \text{ and } (v/2, -v/2, -\sqrt{2}u_2/3, -2u_2/3, u_3), \tag{33}$$

for the $Z$ and $Z_X$ gauge bosons. One exact mass eigenstate is $A_{12} = [2Im(S_1) - \sqrt{2}Im(S_2)]/\sqrt{3}$ with mass given by

$$m_{A_{12}}^2 = -6m_{12}u_2. \tag{34}$$

Assuming that $v^2 << u_{2,3}^2$, the other two mass eigenstates are $A \simeq -Im(\phi_1^0) + Im(\phi_2^0)$ and $A_S \simeq [u_3 Im(S_1) + \sqrt{2}u_3 Im(S_2) + \sqrt{2}u_2 Im(S_3)]/\sqrt{u_2^2 + 3u_3^2/2}$ with masses given by

$$m_A^2 \simeq (4f^2 - 2g_X^2) u_3^2 + \frac{4}{3} g_X^2 u_2^2, \tag{35}$$

$$m_{A_S}^2 \simeq -A_h h \left(\frac{3u_3}{\sqrt{2}} + \frac{\sqrt{2}u_2^2}{u_3}\right), \tag{36}$$

respectively. The charged scalar $H^\pm = (-\phi_1^\pm + \phi_2^\pm)/\sqrt{2}$ has a mass given by

$$m_{H^\pm}^2 = (4f^2 - 2g_X^2) u_3^2 + \frac{4}{3} g_X^2 u_2^2 + (g_2^2 - 2f^2) v^2. \tag{37}$$

## 5. Physical scalars and pseudoscalars

In the MSSM without radiative corrections,

$$m_{H^\pm}^2 = m_A^2 + m_W^2, \tag{38}$$

$$m_{h,H}^2 = \frac{1}{2}\left(m_A^2 + m_Z^2 \mp \sqrt{(m_A^2 + m_Z^2)^2 - 4m_Z^2 m_A^2 \cos^2 2\beta}\right), \tag{39}$$

where $\tan\beta = v_2/v_1$. For $v_1 = v_2$ as in this model, $m_h$ would be zero. There is of course the important radiative correction from Eq. (14), but that alone will not reach 125 GeV. Hence the MSSM requires both large $\tan\beta$ and large radiative correction, but a significant tension remains in accommodating all data. In this model, as Eq. (23) shows, $m_h^2 \simeq (g_X^2 + 2f^2 + \lambda_2)v^2$, where $v = 123$ GeV. This is a very interesting and important result, allowing the Higgs boson mass to be determined by the gauge $U(1)_X$ coupling $g_X$ in addition to the Yukawa coupling $f$ which

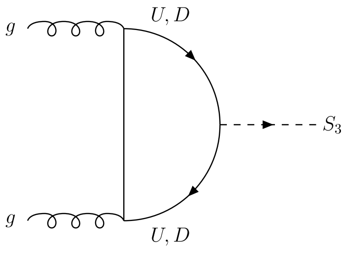


Fig. 1. One-loop production of $S_3$ by gluon fusion.

replaces the $\mu$ parameter, i.e. $\mu = fu_3$. There is no tension between $m_h = 125$ GeV and the superparticle mass spectrum. Since $\lambda_2 \simeq 0.25$ for $\tilde{m}_t \simeq 1$ TeV, we have the important constraint

$$\sqrt{g_X^2 + 2f^2} \simeq 0.885. \tag{40}$$

For illustration, we have already chosen $g_X = 0.53$. Hence $f = 0.5$ and for $u_3 = 2$ TeV, $fu_3 = 1$ TeV is the value of the $\mu$ parameter of the MSSM. Let us choose $u_2 = 4$ TeV, then $m_{Z_X} = 2.87$ TeV, which is slightly above the present experimental lower bound of 2.85 TeV using $g_X = 0.53$ discussed earlier.

As for the heavy Higgs doublet, the four components $(H^{\pm}, H, A)$ are all degenerate in mass, i.e. $m^2 \simeq (4f^2 - 2g_X^2)u_3^2 + (4/3)g_X^2 u_2^2$ up to $v^2$ corrections. Each mass is then about 2.78 TeV. In more detail, as shown in Eq. (37), $m_{H^{\pm}}^2$ is corrected by $g_2^2 v^2 = m_W^2$ plus a term due to $f$. As shown in Eq. (24), $m_H^2$ is corrected by $(g_1^2 + g_2^2)v^2 = m_Z^2$ plus a term due to $f$ and $\lambda_2$. These are exactly in accordance with Eqs. (38) and (39).

In the $S_{1,2,3}$ sector, the three physical scalars are mixtures of all three $Re(S_i)$ components, whereas the physical pseudoscalar $A_{12}$ has no $Im(S_3)$ component. Since only $S_3$ couples to $UU^c$, $DD^c$, and $\phi_1\phi_2$, a candidate for the 750 GeV diphoton resonance must have an $S_3$ component. It could be one of the three scalars or the pseudoscalar $A_S$, or the other $S_3$ without VEV. In the following, we will consider the last option, specifically a pseudoscalar $\chi$ with a significant component of this other $S_3$. This allows the $\chi UU^c$, $\chi DD^c$ and $\chi\phi_1\phi_2$ couplings to be independent of the masses of $U$, $D$, and the charged higgsino. The other scalars and pseudoscalars are assumed to be much heavier, and yet to be discovered.

## 6. Diphoton excess

In this model, other than the addition of $N^c$ for seesaw neutrino masses, the only new particles are $U, U^c, D, D^c$ and $S_{1,2,3}$, which are exactly the ingredients needed to explain the diphoton excess at the LHC. The allowed $S_3UU^c$ and $S_3DD^c$ couplings enable the one-loop gluon production of $S_3$ in analogy to that of $h$ (see Fig. 1). The one-loop decay of $S_3$ to two photons comes from these couplings as well as $S_3\phi_1\phi_2$ (see Fig. 2). In addition, the direct $S_1S_2S_3$ couplings enable the decay of $S_3$ to other final states, including those of the dark sector, which contribute to its total width. The fact that the exotic $U, U^c, D, D^c$ scalars are leptoquarks is also very useful for understanding [15] other possible LHC flavor anomalies. In a nutshell, a desirable comprehensive picture of possible new physics beyond the standard model is encapsulated by this existing model. In the following, we assume that the pseudoscalar $\chi$ is the 750 GeV particle, and show how its production and decay are consistent with the present data.

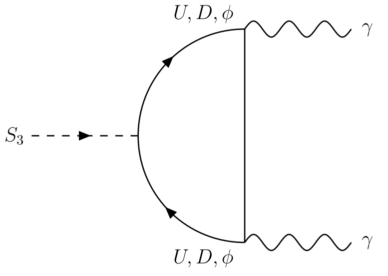


Fig. 2. One-loop decay of $S_3$ to two photons.

The production cross section through gluon fusion is given by

$$\hat{\sigma}(gg \to \chi) = \frac{\pi^2}{8m_\chi}\Gamma(\chi \to gg)\delta(\hat{s} - m_\chi^2). \tag{41}$$

For the LHC at 13 TeV, the diphoton cross section is roughly [16]

$$\sigma(gg \to \chi \to \gamma\gamma) \simeq (100 \text{ pb}) \times (\lambda_g \text{ TeV})^2 \times B(\chi \to \gamma\gamma), \tag{42}$$

where $\lambda_g$ is the effective coupling of $\chi$ to two gluons, normalized by

$$\Gamma(\chi \to gg) = \frac{\lambda_g^2}{8\pi} m_\chi^3. \tag{43}$$

Let the $\chi \bar{Q} Q$ coupling be $f_Q$, where $Q$ is a leptoquark fermion, then

$$\lambda_g = \frac{\alpha_s}{\pi m_\chi} \sum_Q f_Q F(m_Q^2/m_\chi^2), \tag{44}$$

where [17]

$$F(x) = 2\sqrt{x}\left[\arctan\left(\frac{1}{\sqrt{4x-1}}\right)\right]^2, \tag{45}$$

which has the maximum value of $\pi^2/4 = 2.47$ as $x \to 1/4$. Let $f_Q^2/4\pi = 0.21$ and $F(m_Q^2/m_\chi^2) = 2.0$ (i.e. $m_Q = 380$ GeV) for all $Q = U, U, D$, then $\lambda_g = 0.49\ \text{TeV}^{-1}$. For the corresponding

$$\Gamma(\chi \to \gamma\gamma) = \frac{\lambda_\gamma^2}{64\pi} m_\chi^3, \tag{46}$$

the $\phi^\pm$ higgsino contributes as well as $U, D$. However, its mass is roughly $fu_3 = 1$ TeV, so $F(x_\phi) = 0.394$, and

$$\lambda_\gamma = \frac{2\alpha}{\pi m_\chi} \sum_\psi N_\psi Q_\psi^2 f_\psi F(x_\psi), \tag{47}$$

where $\psi = U, U, D, \phi^\pm$ and $N_\psi$ is the number of copies of $\psi$. Using $f_\phi^2/4\pi = 0.21$ as well, $\lambda_\gamma = 0.069\ \text{TeV}^{-1}$ is obtained. We then have $\Gamma(\chi \to \gamma\gamma) = 10$ MeV and $\Gamma(\chi \to gg) = 4.0$ GeV. If $B(\chi \to \gamma\gamma) = 2.5 \times 10^{-4}$, then $\sigma = 6$ fb, and the total width of $\chi$ is 40 GeV, in good agreement with data [1,2].

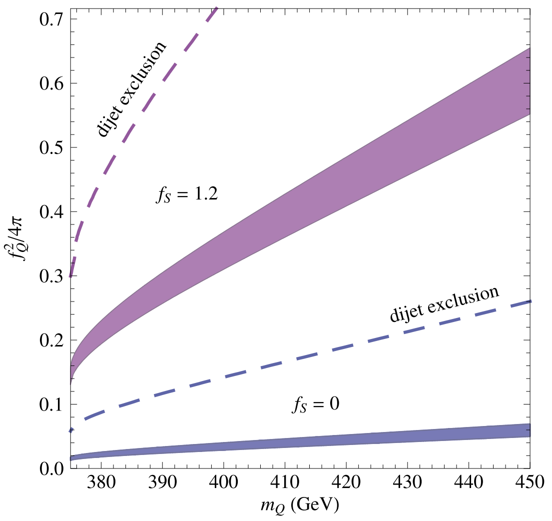


Fig. 3. Allowed region for diphoton cross section of $6.2 \pm 1$ fb.

Note the important fact that we have considered 380 GeV for the mass of the leptoquark fermions. If they are leptoquark scalars, then their mass would be constrained by LHC data to be above 1 TeV or so. As fermions, $Q$ has odd $R$ parity, and must decay into the lightest supersymmetric particle, which is discussed in more detail below. We assume 200 GeV for this particle, hence there is no useful bound on $m_Q$ at present.

As mentioned earlier, there are 2 copies of $S_3$ and 3 copies each of $S_{1,2}$. In addition to the ones with VEVs in their scalar components, there are 5 other superfields. One pair $\tilde{S}_{1,2}$ may form a pseudo-Dirac fermion, and be the lightest particle with odd $R$ parity. It will couple to $\chi$, say with strength $f_S$ which is independent of all other couplings that we have discussed, then the tree-level decay $\chi \to \tilde{S}_1 \tilde{S}_2$ dominates the total width of $\chi$ and is invisible

$$\Gamma(\chi \to \tilde{S}_1 \tilde{S}_2) = \frac{f_S^2}{8\pi}\sqrt{m_\chi^2 - 4m_S^2}. \tag{48}$$

For $m_\chi = 750$ GeV and $m_S = 200$ GeV, we find $\Gamma = 36$ GeV if $f_S = 1.2$. These numbers reinforce our numerical analysis to support the claim that $\chi$ is a possible candidate for the 750 GeV diphoton excess. Note also that $\lambda_g$ and $\lambda_\gamma$ have scalar contributions which we have not considered. Adding them will allow us to reduce the fermion contributions we have assumed and still get the same final results.

If we disregard the decay to dark matter ($f_S = 0$), then the total width of $\chi$ is dominated by $\Gamma(\chi \to gg)$, which is then less than a GeV. Assuming that the cross section for the diphoton resonance is $6.2 \pm 1$ fb [16], we plot the allowed values of $f_Q^2/4\pi$ versus $m_Q$ for both $f_S = 1.2$ which gives a total width of about 40 GeV for $\chi$, and $f_S = 0$ which requires much smaller values of $f_Q^2/4\pi$. Since $\chi$ must also decay into two gluons, we show the diject exclusion upper limits ($\sim 2$ pb) from the 8 TeV data in each case as well. Our choice of the pseudoscalar $\chi$ to be the 750 GeV diphoton resonance is motivated by the necessity of large couplings to $U$, $D$ leptoquark fermions for explaining the large width of about 40 GeV observed by ATLAS. If we take the evidence of CMS that this width is narrow, then as Fig. 3 shows, we can have much smaller couplings and much greater masses for $U, D$. In that case, we can use a physical scalar, with mass-squared matrix given in Eqs. (28) to (32), which is directly associated with the $\mu$ term.

## 7. Scalar neutrino and neutralino sectors

In the neutrino sector, the $2 \times 2$ mass matrix spanning $(\nu, N^c)$ per family is given by the well-known seesaw structure:

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & m_D \\ m_D & m_N \end{pmatrix}, \tag{49}$$

where $m_D$ comes from $v_2$ and $m_N$ from $u_1$. There are two neutral complex scalars with odd $R$ parity per family, i.e. $\tilde{\nu} = (\tilde{\nu}_R + i\tilde{\nu}_I)/\sqrt{2}$ and $\tilde{N}^c = (\tilde{N}^c_R + i\tilde{N}^c_I)/\sqrt{2}$. The $4 \times 4$ mass-squared matrix spanning $(\tilde{\nu}_R, \tilde{\nu}_I, \tilde{N}^c_R, \tilde{N}^c_I)$ is given by

$$\mathcal{M}^2_{\tilde{\nu},\tilde{N}^c} = \begin{pmatrix} m^2_{\tilde{\nu}} & 0 & A_D m_D & 0 \\ 0 & m^2_{\tilde{\nu}} & 0 & -A_D m_D \\ A_D m_D & 0 & m^2_{\tilde{N}^c} + A_N m_N & 0 \\ 0 & -A_D m_D & 0 & m^2_{\tilde{N}^c} - A_N m_N \end{pmatrix}. \tag{50}$$

In the MSSM, $\tilde{\nu}$ is ruled out as a dark-matter candidate because it interacts elastically with nuclei through the $Z$ boson. Here, the $A_N$ term allows a mass splitting between the real and imaginary parts of the scalar fields, and avoids this elastic-scattering constraint by virtue of kinematics. However, we still assume their masses to be heavier than that of $\tilde{S}_{1,2}$, discussed in the previous section.

In the neutralino sector, in addition to the $4 \times 4$ mass matrix of the MSSM spanning $(\tilde{B}, \tilde{W}_3, \tilde{\phi}^0_1, \tilde{\phi}^0_2)$ with the $\mu$ parameter replaced by $fu_3$, i.e.

$$\mathcal{M}_0 = \begin{pmatrix} M_1 & 0 & -g_1 v_1/\sqrt{2} & g_1 v_2/\sqrt{2} \\ 0 & M_2 & g_2 v_1/\sqrt{2} & -g_2 v_2/\sqrt{2} \\ -g_1 v_1/\sqrt{2} & g_2 v_1/\sqrt{2} & 0 & -fu_3 \\ g_1 v_2/\sqrt{2} & -g_2 v_2/\sqrt{2} & -fu_3 & 0 \end{pmatrix}, \tag{51}$$

there is also the $4 \times 4$ mass matrix spanning $(\tilde{X}, \tilde{S}_3, \tilde{S}_2, \tilde{S}_1)$, i.e.

$$\mathcal{M}_S = \begin{pmatrix} M_X & \sqrt{2} g_X u_3 & -2\sqrt{2} g_X u_2/3 & -\sqrt{2} g_X u_1/3 \\ \sqrt{2} g_X u_3 & 0 & h u_1 & h u_2 \\ -2\sqrt{2} g_X u_2/3 & h u_1 & 0 & h u_3 \\ -\sqrt{2} g_X u_1/3 & h u_2 & h u_3 & 0 \end{pmatrix}. \tag{52}$$

The two are connected through the $4 \times 4$ matrix

$$\mathcal{M}_{0S} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ -g_x v_1/\sqrt{2} & -f v_2 & 0 & 0 \\ -g_X v_2/\sqrt{2} & -f v_1 & 0 & 0 \end{pmatrix}. \tag{53}$$

These neutral fermions are odd under $R$ parity and the lightest could in principle be a dark-matter candidate. To avoid the stringent bounds on dark matter with the MSSM alone, we assume again that all these particles are heavier than $\tilde{S}_{1,2}$, as the dark matter discussed in the previous section.

## 8. Dark matter

The $5 \times 5$ mass matrix spanning the 5 singlet fermions $(\tilde{S}_1, \tilde{S}_2, \tilde{S}_1, \tilde{S}_2, \tilde{S}_3)$, corresponding to superfields with zero VEV for their scalar components, is given by

$$\mathcal{M}_{\tilde{S}} = \begin{pmatrix} 0 & m_0 & 0 & 0 & m_{13} \\ m_0 & 0 & 0 & 0 & m_{23} \\ 0 & 0 & 0 & M_3 & M_2 \\ 0 & 0 & M_3 & 0 & M_1 \\ m_{13} & m_{23} & M_2 & M_1 & 0 \end{pmatrix}. \tag{54}$$

Note that the $4 \times 4$ submatrix spanning $(\tilde{S}_1, \tilde{S}_2, \tilde{S}_1, \tilde{S}_2)$ has been diagonalized to form two Dirac fermions. We can choose $m_0$ to be small, say 200 GeV, and $M_{1,2,3}$ to be large, of order TeV. However, because of the mixing terms $m_{13}, m_{23}$, the light Dirac fermion gets split into two Majorana fermions, so it should be called a pseudo-Dirac fermion.

The dark matter with odd $R$ parity is the lighter of the two Majorana fermions, call it $\tilde{S}$, contained in the pseudo-Dirac fermion formed out of $\tilde{S}_{1,2}$ as discussed in Sec. 6. It couples to the $Z_X$ gauge boson, but in the nonrelativistic limit, its elastic scattering cross section with nuclei through $Z_X$ vanishes because it is Majorana. It also does not couple directly to the Higgs boson $h$, so its direct detection at underground search experiments is very much suppressed. However, it does couple to $A_S$ which couples also to quarks through the very small mixing of $A_S$ with $A$. This is further suppressed because it contributes only to the spin-dependent cross section. To obtain a spin-independent cross section at tree level, the constraint of Eqs. (17) to (19) have to be relaxed so that $h$ mixes with $S_{1,2,3}$.

Let the coupling of $h$ to $\tilde{S}\tilde{S}$ be $\epsilon$, then the effective interaction for elastic scattering of $\tilde{S}$ with nuclei through $h$ is given by

$$\mathcal{L}_{eff} = \frac{\epsilon f_q}{m_h^2} \bar{\tilde{S}}\tilde{S}\bar{q}q, \tag{55}$$

where $f_q = m_q/2v = m_q/(246\ \text{GeV})$. The spin-independent direct-detection cross section per nucleon is given by

$$\sigma^{SI} = \frac{4\mu_{DM}^2}{\pi A^2}[\lambda_p Z + (A - Z)\lambda_n]^2, \tag{56}$$

where $\mu_{DM} = m_{DM} M_A/(m_{DM} + M_A)$ is the reduced mass of the dark matter. Using [18]

$$\lambda_N = \left[\sum_{u,d,s} f_q^N + \frac{2}{27}\left(1 - \sum_{u,d,s} f_q^N\right)\right] \frac{\epsilon m_N}{(246\ \text{GeV}) m_h^2}, \tag{57}$$

with [19]

$$f_u^p = 0.023, \quad f_d^p = 0.032, \quad f_s^p = 0.020, \tag{58}$$

$$f_u^n = 0.017, \quad f_d^n = 0.041, \quad f_s^n = 0.020, \tag{59}$$

we find $\lambda_p \simeq 3.50 \times 10^{-8}\ \text{GeV}^{-2}$, and $\lambda_n \simeq 3.57 \times 10^{-8}\ \text{GeV}^{-2}$. Using $A = 131$, $Z = 54$, and $M_A = 130.9$ atomic mass units for the LUX experiment [20], and $m_{DM} = 200$ GeV, we find for the upper limit of $\sigma^{SI} < 1.5 \times 10^{-45}\ \text{cm}^2$, the bound $\epsilon < 6.5 \times 10^{-4}$.

We have already invoked the $\chi\tilde{S}_1\tilde{S}_2$ coupling to obtain a large invisible width for $\chi$. Consider now the fermion counterpart of $\chi$, call it $\tilde{S}'$, and the scalar counterparts of $\tilde{S}_{1,2}$, then the couplings

$\tilde{S}'\tilde{S}_1 S_2$ and $\tilde{S}'\tilde{S}_2 S_1$ are also $f_S = 1.2$. Suppose one linear combination of $S_{1,2}$, call it $\zeta$, is lighter than 200 GeV, then the thermal relic abundance of dark matter is determined by the annihilation $\tilde{S}\tilde{S} \to \zeta\zeta$, with a cross section times relative velocity given by

$$\sigma \times v_{rel} = \frac{f_\zeta^4 m_{S'}^2 \sqrt{1 - m_\zeta^2/m_S^2}}{16\pi (m_{S'}^2 + m_S^2 - m_\zeta^2)^2}. \tag{60}$$

Setting this equal to the optimal value [21] of $2.2 \times 10^{-26}$ cm$^3$/s, we find $f_\zeta \simeq 0.62$ for $m_{S'} = 1$ TeV, $m_S = 200$ GeV, and $m_\zeta = 150$ GeV. Note that $\zeta$ stays in thermal equilibrium through its interaction with $h$ from a term in $V_D$. It is also very difficult to be produced at the LHC, because it is an SM singlet, so its mass of 150 GeV is allowed.

## 9. Conclusion

The utilitarian supersymmetric $U(1)_X$ gauge extension of the Standard Model of particle interactions proposed 14 years ago [4] allows for two classes of anomaly-free models which have no $\mu$ term and conserve baryon number and lepton number automatically. A simple version [7] with leptoquark superfields is especially interesting because of existing LHC flavor anomalies.

The new $Z_X$ gauge boson of this model has specified couplings to quarks and leptons which are distinct from other gauge extensions and may be tested at the LHC. On the other hand, a hint may already be discovered with the recent announcements by ATLAS and CMS of a diphoton excess at around 750 GeV. It may well be the revelation of the singlet scalar (or pseudoscalar) $S_3$ predicted by this model which also predicts that there should be singlet leptoquarks and other particles that $S_3$ must couple to. Consequently, gluon fusion will produce $S_3$ which will then decay to two photons together with other particles, including those of the dark sector. This scenario explains the observed diphoton excess, all within the context of the original model, and not an invention after the fact.

Since $S_3$ couples to leptoquarks, the $S_3 \to l_i^+ l_j^-$ decay must occur at some level. As such, $S_3 \to e^+\mu^-$ would be a very distinct signature at the LHC. Its branching fraction depends on unknown Yukawa couplings which need not be very small. Similarly, the $S_3$ couplings to $\phi_1\phi_2$ as well as leptoquarks imply decays to $ZZ$ and $Z\gamma$ with rates comparable to $\gamma\gamma$.

An important byproduct of this study is the discovery of relaxed supersymmetric constraints on the Higgs boson's mass of 125 GeV. It is now given by Eq. (23), i.e. $m_h^2 \simeq (g_X^2 + 2f^2 + \lambda_2)v^2$, which allows it to be free of the tension encountered in the MSSM. This prediction is independent of whether the diphoton excess is confirmed or not.

Most importantly, since $S_3$ replaces the $\mu$ parameter, its association with the 750 GeV excess implies the existence of supersymmetry. If confirmed and supported by subsequent data, it may even be considered in retrospect as the first evidence for the long-sought existence of supersymmetry.

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